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Spin Pumping and Shot Noise in Ferrimagnets: Bridging Ferro- and Antiferromagnets

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A combination of novel technological and fundamental physics prospects has sparked a huge interest in pure spin transport in magnets, starting with ferromagnets and spreading to antiferro and ferrimagnets. We present a theoretical study of spin transport across a ferrimagnet nonmagnetic conductor interface, when a magnetic eigenmode is driven into a coherent state. The obtained spin current expression includes intra as well as cross sublattice terms, both of which are essential for a quantitative understanding of spin pumping. The dc current is found to be sensitive to the asymmetry in interfacial coupling between the two sublattice magnetizations and the mobile electrons, especially for antiferromagnets. We further find that the concomitant shot noise provides a useful tool for probing the quasiparticle spin and interfacial coupling.

Introduction.—The quest for energy-efficient information technology has driven scientists to examine unconventional means of data transmission and processing. Pure spin current transport in magnetic insulators has emerged as one of the most promising candidates [1–4]. Heterostructures composed of an insulating magnet and a nonmagnetic conductor (N) enable the conversion of the magnonic spin current in the former to the electronic in the latter, thereby allowing for their integration with conventional electronics. In conjunction with the technological pull, these low-dissipation systems have provided a fertile playground for fundamental physics [5–7].

Commencing the exploration with ferromagnets (FMs), the focus in recent years has been shifting towards antiferromagnets (AFMs) [8–10] due to their technological advantages [11]. While a qualitative understanding of some aspects of AFMs, such as spin pumping [12,13], has been borrowed without much change from FMs, the leading-order effects in several other phenomena, such as spin transfer torque [13] and magnetization dynamics [8], bear major qualitative differences. Thus, several phenomena, already known for FMs, are now being generalized for AFMs [14].

Although ferrimagnets ($\mathcal{F}$'s) have been the subject of comparatively fewer works [7,15,16], their high potential is undoubted. The additional complexity of their magnetic structure comes hand in hand with broader possibilities and still newer phenomena. The spin Seebeck effect [17–19] in an $\mathcal{F}$ with a magnetic compensation temperature has unveiled rich physics due to the interplay between the opposite spin excitations in the magnet [16]. Further studies have asserted an important role of the interfacial coupling between the magnet and the conductor [20]. While yttrium iron garnet is a ferrimagnet and has been the subject of several studies [1,3,4,21–23], it is often treated as a ferromagnet on the grounds that only the low-energy magnons are important [24].

In this Letter, we evaluate the spin pumping current (I_{sz}) and the concomitant spin current shot noise [$S(\Omega)$] in a $\mathcal{F} - N$ bilayer [Fig. 1(a)], when one of the $\mathcal{F}$ eigenmodes is driven into a coherent sate. A two-sublattice model with easy-axis anisotropy and a collinear ground state is employed. Our model continuously encompasses systems from ferromagnets to antiferromagnets, thereby allowing analytical results for the full range of materials within a unified description. It further allows arbitrary (disordered) interfaces. In addition to the bulk asymmetry, stemming from inequivalent sublattices, we find a crucial role for the interfacial coupling asymmetry (Fig. 2), consistent with the existing experiments [16,20] and theoretical proposals [25]. Such an asymmetry may occur even in a perfect crystalline interface [Fig. 1(b)] due to the nature of the termination or the different wave function clouds of the electrons

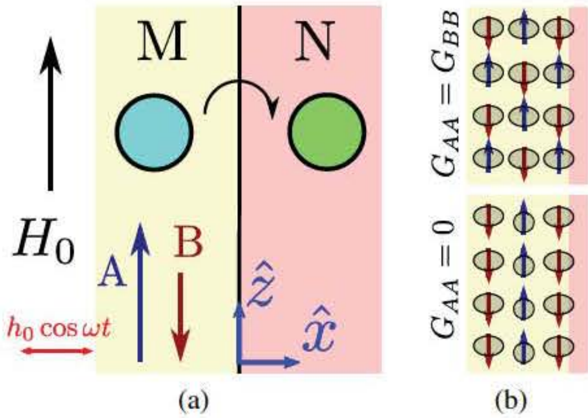

FIG. 1. (a) Schematic of the magnet (M) nonmagnetic conductor (N) heterostructure under investigation. Equilibrium magnetization for sublattices A (blue) and B (red) point along $\hat{z}$ and $-\hat{z}$, respectively. An eigenmode in M is driven coherently and injects z polarized spin current into N. (b) Schematics of possible interface microstructures. Shaded regions around each spin represent the wave function cloud of the localized electrons composing the spin. Our model encompasses compensated as well as uncompensated interfaces including lattice disorder.

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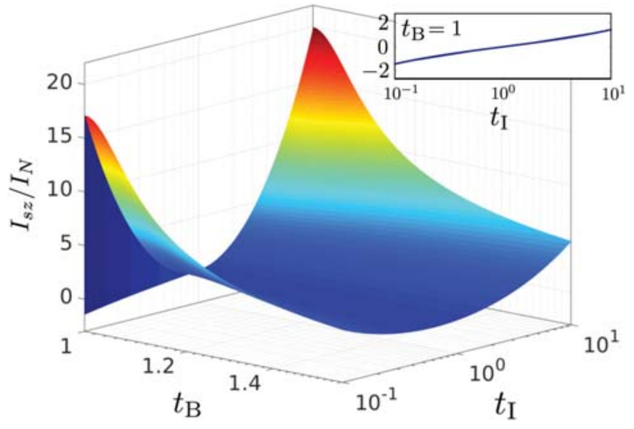

FIG. 2. Normalized spin current vs bulk ($t_B = M_{A0}/M_{B0}$) and interfacial ($t_I = \Gamma_{AA}/\Gamma_{BB}$) asymmetries for the lower frequency uniform mode in a coherent state. All other bulk parameters are kept constant, no external magnetic field is applied, and $I_N = 2\hbar|\chi|^2\omega_q\alpha_{AB}$. The spin current for $t_B = 1$ (also depicted in the inset for clarity) is small due to the spin zero quasiparticles in symmetric AFMs, and it abruptly increases with a small bulk symmetry breaking due to quasiparticle transformation into spin $\hbar$ magnons [7]. The different parameter values employed are given in Supplemental Material [26].

constituting the localized spins in the two sublattices. Spin transport in AFM-N bilayers is found to be particularly sensitive to the interfacial asymmetry, with the spin current nearly vanishing for symmetrical coupling of the two sublattices with N corresponding to the case of a compensated interface (Fig. 2).

A key result of our work is the following semiclassical expression for the spin current injected into N [27]:

$$\frac{e}{\hbar} I_{sz} = \sum_{i,j=\{A,B\}} G_{ij}(\hat{\boldsymbol{m}}_i \times \dot{\hat{\boldsymbol{m}}}_j)_z = \sum_{i,j=\{m,n\}} G_{ij}(\boldsymbol{i} \times \dot{\boldsymbol{j}})_z, \quad (1)$$

where $\hat{\boldsymbol{m}}_{A(B)}$ is the unit vector along sublattice A (B) magnetization, $\boldsymbol{m} = [\hat{\boldsymbol{m}}_A + \hat{\boldsymbol{m}}_B]/2$, $\boldsymbol{n} = [\hat{\boldsymbol{m}}_A - \hat{\boldsymbol{m}}_B]/2$, $G_{mm} = G_{AA} + G_{BB} + 2G_{AB}$, $G_{nn} = G_{AA} + G_{BB} - 2G_{AB}$, and $G_{mn} = G_{nm} = G_{AA} - G_{BB}$. Employing $G_{AB} = G_{BA} = \sqrt{G_{AA}G_{BB}}$, which is derived, along with the expressions for G_{AA} and G_{BB}, in the subsequent discussion below, we further obtain $G_{mm} = (\sqrt{G_{AA}} + \sqrt{G_{BB}})^2$ and $G_{nn} = (\sqrt{G_{AA}} - \sqrt{G_{BB}})^2$. Our result [Eq. (1)] for the injected spin current adds upon the existing understanding of spin pumping via AFMs [13] by (i) providing analytic and intuitive expressions for the conductances, (ii) incorporating the cross terms characterized by G_{AB} and G_{mn}, (iii) deriving the relation $G_{AB} = \sqrt{G_{AA}G_{BB}}$ based upon a microscopic interfacial exchange coupling model, (iv) accommodating compensated ($G_{AA} = G_{BB}$) as well as uncompensated interfaces, and (v) allowing for interfacial disorder. As detailed in Supplemental Material [26], the spin pumping expression given in Ref. [13] is recovered from Eq. (1) by substituting $G_{AB} = G_{BA} = 0$ and $G_{AA} = G_{BB}$ and yields results qualitatively different from what is reported herein [26]. This difference in results stems from the assumption made in Ref. [13] that $\hat{\boldsymbol{m}}_A$ and $\hat{\boldsymbol{m}}_B$ are independent variables, which is equivalent to setting $G_{AB} = G_{BA} = 0$ implicitly. $\hat{\boldsymbol{m}}_A$ and $\hat{\boldsymbol{m}}_B$ are coupled via intersublattice exchange and hence cannot be treated as independent when considering system dynamics.

We define the dynamical spin correction factor S_D via the relation $S_D \equiv \lim_{T\to 0} S(0)/2\hbar I_{sz}$, where T is the temperature and $S(0)$ is the low-frequency spin current shot noise. When the effect of either the dipolar interaction [29] or the sublattice coupling on the eigenmode under consideration can be disregarded, $S_D\hbar$ coincides with the spin of the eigenmode. In other words, when a full four-dimensional (4D) Bogoliubov transform [7] is required to obtain the relevant eigenmode, S_D is a property of the entire heterostructure and depends upon the bulk as well as the interface. Thus, shot noise offers a useful experimental probe of the interfacial properties as discussed below.

Model.—The model we study consists of a two-sublattice magnet coupled via an interfacial exchange interaction to a nonmagnetic conductor [Fig. 1(a)]. We assume $M_{A0} \geq M_{B0}$ with the respective sublattice saturation magnetizations M_{A0} and M_{B0}. The bulk of the magnet is characterized by a classical free energy density which is then quantized, using the Holstein-Primakoff transformations [30–32], to yield the magnetic contribution to the quantum Hamiltonian $\tilde{\mathcal{H}}_M$ in terms of the magnon ladder operators.

We consider Zeeman (H_Z), easy-axis anisotropy ($H_{\rm an}$), exchange ($H_{\rm ex}$), and dipolar interaction ($H_{\rm dip}$) (see Ref. [29]) in the magnetic free energy density written in terms of the A and B sublattice magnetizations $\boldsymbol{M}_A(\boldsymbol{r})$ and $\boldsymbol{M}_B(\boldsymbol{r})$. With an applied magnetic field $H_0\hat{\boldsymbol{z}}$ and μ_0 the permeability of free space, the Zeeman energy density reads $H_Z = -\mu_0 H_0(M_{Az} + M_{Bz})$. The easy-axis anisotropy is parametrized in terms of the constants K_{uA} and K_{uB} as $H_{\rm an} = -K_{uA}M_{Az}^2 - K_{uB}M_{Bz}^2$ [32]. The exchange energy density is expressed in terms of the constants $\mathcal{J}_A$, $\mathcal{J}_B$, $\mathcal{J}_{AB}$, and $\mathcal{J}$ [32]: $H_{\rm ex} = \sum_{x_i=x,y,z}[\mathcal{J}_A(\partial\boldsymbol{M}_A/\partial x_i)\cdot(\partial\boldsymbol{M}_A/\partial x_i) + \mathcal{J}_B(\partial\boldsymbol{M}_B/\partial x_i)\cdot(\partial\boldsymbol{M}_B/\partial x_i) + \mathcal{J}_{AB}(\partial\boldsymbol{M}_A/\partial x_i)\cdot(\partial\boldsymbol{M}_B/\partial x_i)] + \mathcal{J}\boldsymbol{M}_A\cdot\boldsymbol{M}_B$. The dipolar interaction energy density is obtained in terms of the demagnetization field $\boldsymbol{H}_m$ that obeys Maxwell's equations in the magnetostatic approximation: $H_{\rm dip} = -(1/2)\mu_0\boldsymbol{H}_m\cdot(\boldsymbol{M}_A + \boldsymbol{M}_B)$ [7,31,32]. Quantizing the magnetization fields and employing the Holstein-Primakoff transformation, we obtain the quantum Hamiltonian for the magnet:

$$\tilde{\mathcal{H}}_M = \sum_{\boldsymbol{q}}\left(\frac{A_{\boldsymbol{q}}}{2}\tilde{a}_{\boldsymbol{q}}^\dagger\tilde{a}_{\boldsymbol{q}} + \frac{B_{\boldsymbol{q}}}{2}\tilde{b}_{\boldsymbol{q}}^\dagger\tilde{b}_{\boldsymbol{q}} + C_{\boldsymbol{q}}\tilde{a}_{\boldsymbol{q}}\tilde{b}_{\ \boldsymbol{q}} + D_{\boldsymbol{q}}\tilde{a}_{\boldsymbol{q}}\tilde{a}_{\ \boldsymbol{q}} + E_{\boldsymbol{q}}\tilde{b}_{\boldsymbol{q}}\tilde{b}_{\ \boldsymbol{q}} + F_{\boldsymbol{q}}\tilde{a}_{\boldsymbol{q}}\tilde{b}_{\boldsymbol{q}}^\dagger\right) + \text{H.c.}, \quad (2)$$

where $\tilde{a}_{\boldsymbol{q}}$ and $\tilde{b}_{\boldsymbol{q}}$ are, respectively, sublattice A and B magnon annihilation operators corresponding to wave vector $\boldsymbol{q}$. Relegating the detailed expressions for the coefficients $A_{\boldsymbol{q}}, B_{\boldsymbol{q}}, \ldots$ to Supplemental Material [26], we note that $C_{\boldsymbol{q}}$ is dominated by the intersublattice exchange while $D_{\boldsymbol{q}}$, $E_{\boldsymbol{q}}$, and $F_{\boldsymbol{q}}$ result entirely from dipolar interaction. The magnetic Hamiltonian is diagonalized via a 4D Bogoliubov transform to new operators [7] $\tilde{\alpha}_{\boldsymbol{q}} = u_{l\boldsymbol{q}}\tilde{a}_{\boldsymbol{q}} + v_{l\boldsymbol{q}}\tilde{b}^{\dagger}{}_{\boldsymbol{q}} + w_{l\boldsymbol{q}}\tilde{a}^{\dagger}{}_{\boldsymbol{q}} + x_{l\boldsymbol{q}}\tilde{b}_{\boldsymbol{q}}$ and similar for $\tilde{\beta}_{\boldsymbol{q}}$: $\tilde{\mathcal{H}}_M = \sum_{\boldsymbol{q}}\hbar\omega_{l\boldsymbol{q}}\tilde{\alpha}^{\dagger}_{\boldsymbol{q}}\tilde{\alpha}_{\boldsymbol{q}} + \hbar\omega_{u\boldsymbol{q}}\tilde{\beta}^{\dagger}_{\boldsymbol{q}}\tilde{\beta}_{\boldsymbol{q}}$. The subscripts l and u refer to lower and upper modes, respectively, thus assigning the lower energy to $\tilde{\alpha}$ modes. The diagonal eigenmodes are dressed magnons with the spin given by $\hbar(|u_{\boldsymbol{q}}|^2 - |v_{\boldsymbol{q}}|^2 + |w_{\boldsymbol{q}}|^2 - |x_{\boldsymbol{q}}|^2)$ [7]. Disregarding dipolar interaction, the eigenmode spin is plus or minus $\hbar$. Incorporating the dipolar contribution, the spin magnitude varies between 0 and greater than $\hbar$ [7].

The nonmagnetic conductor is modeled as a bath of noninteracting electrons: $\tilde{\mathcal{H}}_N = \sum_{\boldsymbol{k},s=\pm}\hbar\omega_{\boldsymbol{k}}\tilde{c}^{\dagger}_{\boldsymbol{k},s}\tilde{c}_{\boldsymbol{k},s}$, where $\tilde{c}_{\boldsymbol{k},s}$ is the annihilation operator corresponding to an electron state with spin $s\hbar/2$ along the z direction and orbital wave function $\psi_{\boldsymbol{k}}(\boldsymbol{r})$. The conductor is coupled to the two sublattices in the magnet via an interfacial exchange interaction parameterized by $\mathcal{J}_{iA}$ and $\mathcal{J}_{iB}$:

$$\tilde{\mathcal{H}}_{\text{int}} = -\frac{1}{\hbar^2}\int_{\mathcal{A}} d^2\varrho \sum_{\mathcal{G}=A,B}[\mathcal{J}_{i\mathcal{G}}\tilde{\boldsymbol{S}}_{\mathcal{G}}(\boldsymbol{\varrho})\cdot\tilde{\boldsymbol{S}}_N(\boldsymbol{\varrho})], \qquad (3)$$

where $\boldsymbol{\varrho}$ is the interfacial position vector, $\mathcal{A}$ is the interfacial area and $\tilde{\boldsymbol{S}}_A$, $\tilde{\boldsymbol{S}}_B$, and $\tilde{\boldsymbol{S}}_N$ represent spin density operators corresponding to the magnetic sublattices A and B and the conductor, respectively. In terms of the eigenmode ladder operators, the interfacial exchange Hamiltonian reduces to [33]

$$\tilde{\mathcal{H}}_{\text{int}} = \hbar\sum_{\boldsymbol{k}_1,\boldsymbol{k}_2,\boldsymbol{q}_1}(\tilde{\mathcal{P}}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1} + \tilde{\mathcal{P}}^{\dagger}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1}), \qquad (4)$$

where $\tilde{\mathcal{P}}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1} \equiv \tilde{c}^{\dagger}_{\boldsymbol{k}_1,+}\tilde{c}_{\boldsymbol{k}_2,}\ (W^A_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1}\tilde{a}_{\boldsymbol{q}_1} + W^B_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1}\tilde{b}^{\dagger}{}_{\boldsymbol{q}_1})$, $\hbar W^{\mathcal{G}}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1} = \mathcal{J}_{i\mathcal{G}}\sqrt{M_{\mathcal{G}0}/2|\gamma_{\mathcal{G}}|\hbar}\int_{\mathcal{A}} d^2\boldsymbol{\varrho}[\psi^*_{\boldsymbol{k}_1}(\boldsymbol{\varrho})\psi_{\boldsymbol{k}_2}(\boldsymbol{\varrho})\phi_{\boldsymbol{q}_1}(\boldsymbol{\varrho})]$ with $\gamma_{\mathcal{G}}$ the typically negative gyromagnetic ratio corresponding to sublattice $\mathcal{G}$ $(= A, B)$, and $\phi_{\boldsymbol{q}_1}(\boldsymbol{r})$ is a wave function of the magnon eigenmode with wave vector $\boldsymbol{q}_1$. Our goal is to examine the spin [34] current and its noise when one of the magnetic eigenmodes is in a coherent state. We may, for example, achieve the $\alpha_{\boldsymbol{q}}$ mode in a coherent state by including a driving term in the Hamiltonian: $\tilde{\mathcal{H}}_{\text{drive}} \sim \cos(\omega_{\boldsymbol{q}}t)(\tilde{\alpha}_{\boldsymbol{q}} + \tilde{\alpha}^{\dagger}_{\boldsymbol{q}})$ [35].

The operator corresponding to the z-polarized spin current injected by M into N is obtained from the interfacial contribution to the time derivative of the total electronic spin $(\tilde{\boldsymbol{S}})$:

$$\tilde{I}_{sz} = \frac{1}{i\hbar}[\tilde{S}_z, \tilde{H}_{\text{int}}] = \hbar\sum_{\boldsymbol{k}_1,\boldsymbol{k}_2,\boldsymbol{q}_1}(-i\tilde{\mathcal{P}}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1} + i\tilde{\mathcal{P}}^{\dagger}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}_1}). \qquad (5)$$

The above definition captures the spin pumping contribution to the current injected into N and disregards the effect of interfacial spin-orbit coupling [36]. The power spectral density of spin current noise $S(\Omega)$ is given by [37] $S(\Omega) = \int_{\infty}^{\infty}\lim_{\tau_0\to\infty}(1/2\tau_0)\int_{\tau_0}^{\tau_0}\langle\tilde{\delta I}_{sz}(\tau)\tilde{\delta I}_{sz}(\tau-t) + \tilde{\delta I}_{sz}(\tau-t)\times\tilde{\delta I}_{sz}(\tau)\rangle d\tau e^{i\Omega t}dt$, where $\langle\rangle$ denotes the expectation value and $\tilde{\delta I}_{sz} = \tilde{I}_{sz} - \langle\tilde{I}_{sz}\rangle$ is the spin current fluctuation operator.

Results and discussion.—The spin current I_{sz} in the steady state is obtained by evaluating the expectation value of the spin current operator $\tilde{I}_{sz}$ [Eq. (5)] assuming a magnetic mode, e.g., $\alpha_{\boldsymbol{q}}$, in a coherent state so that $\tilde{\alpha}_{\boldsymbol{q}}$ may be substituted by a c number χ [38]:

$$I_{sz} = 2\hbar|\chi|^2[\Gamma_{AA}(|u|^2 - |w|^2) + \Gamma_{BB}(|v|^2 - |x|^2) - 2\Gamma_{AB}\text{Re}(u^*v - wx^*)], \qquad (6)$$

where u, v, w, and x correspond to the excited eigenmode, $\Gamma_{ij} = \pi\sum_{\boldsymbol{k}_1,\boldsymbol{k}_2}W^i_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}}(W^j_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}})^*(n_{\boldsymbol{k}_2} - n_{\boldsymbol{k}_1})\delta(\omega_{\boldsymbol{k}_1} - \omega_{\boldsymbol{k}_2} - \omega_{\boldsymbol{q}})$ [39], with $i, j = \{A, B\}$ and $n_{\boldsymbol{k}}$ representing the occupancy of the corresponding electron state given by the Fermi-Dirac distribution. Assuming (i) $W^{\mathcal{G}}_{\boldsymbol{k}_1\boldsymbol{k}_2\boldsymbol{q}}$ depends only on the electron chemical potential μ in N such that it may be substituted by $W^{\mathcal{G}}_{\mu}$ and (ii) the electron density of states around the chemical potential $g(\mu)$ is essentially constant, we obtain the simplified relations: $\Gamma_{ij} = \alpha_{ij}\omega_{\boldsymbol{q}}$. Here, $\alpha_{ij} = \pi\hbar^2 W^i_{\mu}(W^j_{\mu})^* V^2_N g^2(\mu)$, with V_N the volume of N. This also entails $\alpha_{AB} = \alpha_{BA} = \sqrt{\alpha_{AA}\alpha_{BB}}$. Since the classical dynamics of a harmonic mode is captured by the system being in a coherent state [40], the spin current evaluated within our quantum model [Eq. (6)] must be identical to the semiclassical expression expected from the spin pumping theory [12] generalized to a two-sublattice system. As detailed in Supplemental Material [26], we evaluate the semiclassical expression given by Eq. (1) for such a coherent state. The result thus obtained is identical to Eq. (6), provided we identify $G_{ij} = (\alpha_{ij}e/\hbar)\times\sqrt{M_{i0}M_{j0}/|\gamma_i||\gamma_j|}$. Since $\alpha_{AB} = \sqrt{\alpha_{AA}\alpha_{BB}}$, we obtain $G_{AB} = G_{BA} = \sqrt{G_{AA}G_{BB}}$ [41]. These relations along with Eq. (1) constitute one of the main results of this Letter.

In order to gain an understanding of the qualitative physics at play, we examine the injected spin current normalized by $I_N = 2\hbar|\chi|^2\omega_{\boldsymbol{q}}\alpha_{AB}$ around the anticrossing point in the dispersion of a ferrimagnet (Fig. 3) for symmetric interfacial coupling $(\Gamma_{AA} = \Gamma_{BB})$. Because of the dipolar interaction [29], the dressed magnon spin smoothly changes between plus and minus $\hbar$ resulting in a similar smooth transition in the spin current [7]. Figure 2 depicts the normalized spin current injected by the lower-frequency uniform mode $(\boldsymbol{q} = \boldsymbol{0})$ with respect to asymmetries in the bulk t_B $(= M_{A0}/M_{B0})$ and the interface t_I $(= \Gamma_{AA}/\Gamma_{BB})$. For simplicity, we keep all other bulk parameters constant and assume the applied field to vanish. For the case of a perfect AFM $(t_B = 1)$ [42], we find a small

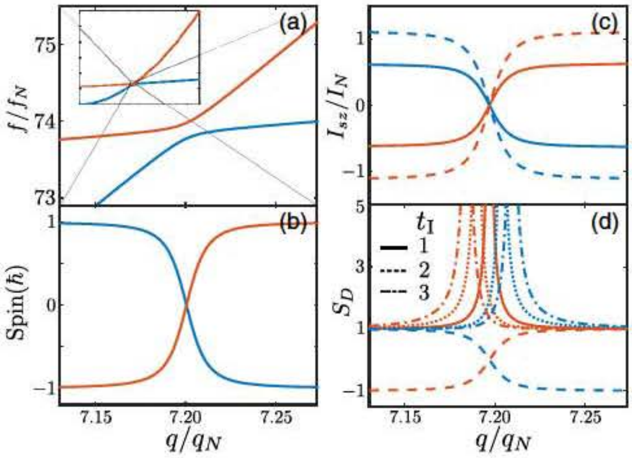

FIG. 3. (a) Dispersion, (b) quasiparticle spin, (c) spin current injected into N, and (d) dynamical spin correction factor vs wave number (along the x direction) around the anticrossing point in a ferrimagnet. $2\pi f_N = |\gamma_A|\mu_0 M_{A0}$ and $f_l(q_N) = 2f_l(0)$ define the normalizations f_N and q_N with $f_l(q)$ the lower dispersion band. $I_N = 2\hbar|\chi|^2\omega_q\alpha_{AB}$ and $t_I \equiv \Gamma_{AA}/\Gamma_{BB} = 1$, unless stated otherwise. The inset in (a) depicts the full dispersion diagram. Dashed lines in (c) depict the spin current I'_{sz} disregarding the cross sublattice terms. Dashed lines in (d) depict the quasiparticle spin, once again, to help comparison. The parameters employed in the plot are given in Supplemental Material [26].

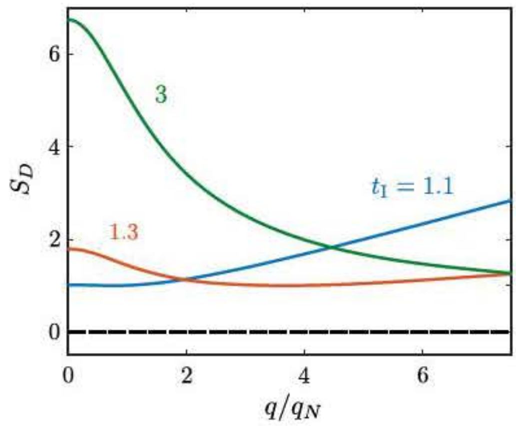

FIG. 4. Dynamical spin correction factor S_D vs wave number (along the x direction) for a symmetrical AFM. The dashed line depicts the zero spin of the magnetic quasiparticles. $f_l(q_N) = 2f_l(0)$ defines the normalization q_N with $f_l(q)$ the lower dispersion band. The parameters employed in the plot are given in Supplemental Material [26].

current with varying t_I that vanishes at $t_I = 1$ (inset in Fig. 2). The small magnitude of the current is attributed to the dipolar interaction-mediated spin-zero magnons in perfect AFMs. The spin current has much larger values when $t_B \neq 0$, since the dressed magnons acquire spin $\hbar$ with a small bulk symmetry breaking [7]. The spin current in this case is highly sensitive to t_I. This sensitivity is particularly pronounced for AFMs, for which the bulk symmetry can also be broken by an applied magnetic field.

The shot noise accompanying the dc spin current injected into N is evaluated for a temperature T:

$$S(\Omega) = 2\hbar|\chi|^2[\alpha_{AA}(|u|^2 + |w|^2) + \alpha_{BB}(|v|^2 + |x|^2) - 2\alpha_{AB}\mathrm{Re}(u^* v + w x^*)][F(\Omega) + F(-\Omega)], \tag{7}$$

where $F(\Omega) \equiv \hbar(\Omega + \omega_q)\coth(\hbar[\Omega + \omega_q]/[2k_B T])$ with k_B the Boltzmann constant. $F(\Omega) \to \hbar|\Omega + \omega_q|$ when $T \to 0$. When the dipolar interaction effect is neglected, i.e., $w, x \to 0$, $\lim_{T\to 0} S(0) \to 2\hbar I_{sz}$ [Eqs. (6) and (7)] such that the dynamical spin correction factor $S_D \to 1$. And when the mode under consideration is not affected by sublattice B, we have $v, x \to 0$ and $S_D\hbar$ approaches the spin of the squeezed magnon [38]. In the general case, $S_D(\geq 1)$ depends upon the magnetic mode and interfacial interaction as well as the eigenmodes in N and is thus a property of the entire heterostructure. Figure 3(d) depicts S_D for a ferrimagnet around the anticrossing point in its dispersion. $S_D \approx 1$ away from the anticrossing and diverges at some wave number which depends upon the interfacial asymmetry t_I. This divergence results from a vanishing I_{sz}. S_D vs the wave number for a symmetric AFM with varying interfacial asymmetry is depicted in Fig. 4. Thus, a combined knowledge of I_{sz} and S_D may allow us to probe interfacial asymmetries experimentally [43]. Since deviations of S_D from 1 are necessarily accompanied by quasiparticles with spin different from $\hbar$, it also offers an indirect signature of their formation.

In order to simplify expressions, we have employed the approximation $W^{\mathcal{G}}_{k_1 k_2 q} \approx W^{\mathcal{G}}_{\mu}$, which is commonly used in the tunneling Hamiltonian description of spin [37,38,44,45] and charge [46] transport. This approximation provides a reasonable description in the limit of strong scattering in N and a disordered interface. The opposite limit of quasiballistic transport in N and an ideal AFM-N interface has been described numerically [13,25,47] as well as analytically [48]. Our approximation further disregards the dependence of the spin conductances on $\boldsymbol{q}$ [49,50].

Summary.—We have presented a theoretical discussion of spin transport across a magnet–nonmagnetic conductor interface when a magnetic eigenmode is driven to a coherent state. Analytical expressions for the dc spin current, including cross terms which were disregarded in Ref. [13], and spin conductances have been obtained. Our theory takes into account the important role of bulk and interfacial sublattice asymmetries as well as lattice disorder at the interface. The spin current, especially in antiferromagnets, is found to be sensitive to interfacial asymmetry. We have evaluated the spin current shot noise at finite temperatures and shown that it can be employed to gain essential insights into quasiparticle spin and interfacial asymmetry.

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Note added in proof.—Recently, Liu and co-workers reported [51] a first principles calculation of damping in metallic antiferromagnets. Their conclusions are fully consistent with our work and show the important role of cross-sublattice terms.

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[27] We emphasize that this expression is restricted to the z component of spin and may not be employed for the full spin vector. It has been shown that the microscopic matrix elements corresponding to x and y polarized spin transport are, in general, different [28]. This distinction is often not made in the literature.
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